

# Multi high charged scalars in the LHC searches and Majorana neutrino mass generations

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## Abstract

Motivated by the large rate of  $H \rightarrow \gamma\gamma$  at the LHC, we explore a class of models with high dimensional representations of scalars to realize small Majorana neutrino masses at two-loop level without imposing any new symmetry. In these models, multi scalars with the electric charges higher than two are naturally expected, which not only enhance the  $H \rightarrow \gamma\gamma$  rate, but provide more searching grounds at the LHC.

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*Introduction*— The Higgs doublet ( $H$ ) not only breaks the electroweak gauge symmetries  $SU(2)_L \times U(1)_Y$  to  $U(1)_Q$  but also gives masses to the charged fermions in the standard model (SM). The main goals of the Large Hadron Collider (LHC) are the search of the Higgs boson ( $H$ ) and the investigation of the electroweak symmetry breaking mechanism. Recently, a boson with its mass around 125 GeV has been discovered in both ATLAS [1] and CMS [2] collaborations, which is most likely the Higgs particle as its properties are consistent with the SM predictions except the large production rate of  $H \rightarrow \gamma\gamma$ . Currently, the excess in both experiments is around  $2\sigma$ . If the deviation is sustained, it is clearly a call for new physics. One of natural mechanisms is to include new charged particles in the SM [3], which would enhance the decay rate due to the new charged loop contributions.

On the other hand, neutrino oscillations observed by the solar, atmospheric, and reactor neutrino experiments reveal that neutrinos are massive but tiny and mix with each other, implying the need of the SM extension. Theoretical studies on model buildings to understand the small neutrino masses are enormously varied if it is allowed to freely invent new particles, scalars and/or fermions, beyond the SM ones. Without a new fermion, the only possibility of having non-zero neutrino masses is that neutrinos are Majorana fermions with the accidental global lepton symmetry in the SM either explicitly or spontaneously broken. Apart from the seesaw mechanism [4], models with tiny Majorana neutrino masses arising from quantum corrections have been proposed [5, 6]. It is clear that this type of theories is usually incorporated with some extension of the scalar sector in the SM. In particular, singly and doubly charged scalars are required in the radiative neutrino mass generation mechanisms at one and two-loop levels, respectively. As a result, the existence of new charged particles is a generic feature of the radiative neutrino models. These charged scalars in turn would help to resolve the excess of the  $H \rightarrow \gamma\gamma$  rate at the LHC as discussed in the literature [7]. Note that in most of the above radiative models, since only  $SU(2)_L$  singlet scalars are introduced, new physics effects are limited in the lepton sector, whereas those involving hadrons, such as the neutrinoless double beta decays believed as a benchmark of the Majorana nature of neutrinos, do not show up. In Ref. [8], an unconventional neutrino mass generation model with a triplet scalar was proposed, leading to two doubly charged scalars with the neutrinoless double beta decays dominated by the short distance contribution. However, this type of models needs adding some *ad hoc* discrete symmetry to forbid the neutrino mass term at tree level [9], which makes these models less appealing. In this study,

we explore the most general cases with arbitrary  $SU(2)_L$  representations for scalars to avoid the tree level neutrino mass generation without imposing any new symmetry. It turns out that a class of models with high odd-dimensional scalars, which generate neutrino masses at two-loop level, exists naturally. Consequently, multi scalar bosons with the electric charges higher than 2-unit appear.

*Model*— For all non-Higgs like scalars with non-trivial  $SU(2)_L \times U(1)_Y$  quantum numbers [10], there are only three possible renormalizable Yukawa interactions, given by

$$f_{ab}\bar{L}_a^c L_b s, \quad y_{ab}\bar{\ell}_{Ra}^c \ell_{Rb} \Phi, \quad \text{and} \quad g_{ab}\bar{L}_a^c L_b T, \quad (1)$$

where  $L$  ( $\ell_R$ ) stands for the left-handed (right-handed) lepton,  $a$  or  $b$  denotes  $e, \mu$  and  $\tau$ ,  $c$  represents the charged conjugation,  $s$  and  $\Phi$  are  $SU(2)_L$  singlet scalar fields with  $Y = 2$  and  $Y = 4$ , and  $T$  is an  $SU(2)_L$  triplet with the hypercharge  $Y = 2$ , respectively [11]. Note that the triplet scalar  $T$  can generate neutrino masses at tree level after developing its vacuum expectation value (VEV), known as the Type-II seesaw mechanism [12], while  $s$  and  $\Phi$  are used in Zee [5] and Zee-Babu [6] models, respectively.

In Eq. (1), the third Yukawa interaction is the most troublesome as it generates neutrino masses at tree level with an extreme small value of the VEV or Yukawa couplings, which is obviously un-natural. Without introducing the triplet  $T$ , the interactions in Eq. (1) are precisely given by the Zee-Babu model [6], which has been extensively studied in the literature, in particular its phenomenology of the doubly charged scalar at the LHC. In this work, we consider a new class of models by adding a scalar field  $\xi$  with  $Y = 2$  and a non-trivial  $SU(2)_L$  representation  $\mathbf{n}$ . To minimize our models, we disregard the singlet scalar  $s$  and keep the other singlet  $\Phi$  so that only the second Yukawa interaction in Eq. (1) can exist at tree level.

The general scalar potential reads

$$\begin{aligned} V(H, \xi, \Phi^{\pm\pm}) = & -\mu_H^2 |H|^2 + \lambda_H |H|^4 + \mu_\xi^2 |\xi|^2 + \lambda_\xi^\alpha |\xi|_\alpha^4 \\ & + \mu_\Phi^2 |\Phi|^2 + \lambda_\Phi |\Phi|^4 + \lambda_{H\xi}^\beta (|H|^2 |\xi|^2)_\beta \\ & + \lambda_{H\Phi} |H|^2 |\Phi|^2 + \lambda_{\xi\Phi} |\xi|^2 |\Phi|^2 + [\mu \xi \xi \Phi + \text{h.c.}], \end{aligned} \quad (2)$$

where  $\alpha, \beta$  are the short-handed notations denoting the possible invariant terms for higher representations in general. Also notice that all terms in the potential are self-hermitian except the last  $\mu$ -term, which is related to the dynamics of the lepton number breaking

to be discussed later. For the even dimensional representations, *i.e.*  $\mathbf{n} = \mathbf{2}, \mathbf{4}, \mathbf{6} \dots$ , the products  $\xi\xi$  vanish since

$$\begin{aligned}\xi\xi &= \epsilon_{ii'}\epsilon_{jj'}\epsilon_{kk'}\dots\xi_{ijk\dots}\xi_{i'j'k'\dots} \\ &= -\epsilon_{i'i}\epsilon_{j'j}\epsilon_{k'k}\dots\xi_{ijk\dots}\xi_{i'j'k'\dots} = 0,\end{aligned}\tag{3}$$

due to the anti-symmetric matrix of  $\epsilon_{ij}$  ( $i, j = 1, 2$ ). Subsequently, we only need to consider the odd dimensional representations of  $\xi$ , *i.e.*  $\mathbf{n} = \mathbf{3}, \mathbf{5}, \mathbf{7} \dots$ . Since the triplet has been dropped out, the next minimal choice is  $\mathbf{n} = \mathbf{5}$ . From now on, we concentrate on this minimal one, *the quintuplet*, with  $\xi = (\xi^{+++}, \xi^{++}, \xi^+, \xi^0, \xi^-)^T$ . One shall bear in mind that the results can be easily extended to those with higher representations of  $\xi$  [13]. In three cases, there are three and two irreducible terms for  $|\xi|^4$  and  $|H|^2|\xi|^2$  respectively.

We now move to the lepton number ( $L$ ) violation. In general,  $U(1)_L$  can be either global or gauge symmetry. If the lepton number indeed comes from a global symmetry as that in the SM, the spontaneous symmetry breaking will generate a Nambu-Goldstone (NG) boson, usually called Majoron. In this model, the VEV  $\langle \xi^0 \rangle$  breaks  $U(1)_L$  spontaneously [15]. Since  $\xi$  is an  $SU(2)$  multiplet, its corresponding Majoron has a direct coupling to the  $Z$  boson, which is strongly constrained by the LEP measurement of the invisible  $Z$  decay width. To resurrect it, we illustrate here by adding an another  $SU(2)_L$  triplet ( $\mathbf{3}$ ) scalar field  $\Delta = (\Delta^+, \Delta^0, \Delta^-)^T$  with  $Y = 0$  to the model, which results in one additional non-hermitian term  $\Delta\xi H^* H^*$ . It is easy to see that the coexistences of  $\bar{\ell}_R^c \ell_R \Phi$ ,  $\mu \xi \xi \Phi$  and  $\Delta \xi H^* H^*$  break the lepton number explicitly.

*Neutrino masses*— In the Zee-Babu model, neutrino masses are generated from the two-loop diagrams due to the couplings of  $s^\pm s^\pm \Phi^{\mp\mp}$ . whereas in our model without  $s^\pm$ , they can be induced from similar two-loop diagrams with  $\Phi^{\pm\pm}$  coupling to  $W^\mp W^\mp$  though the mixing of  $\xi$  as shown in Fig. 1. In this mechanism, neutrino masses are calculable, given by

$$\begin{aligned}m_{\nu_{ab}} &\simeq \frac{g^4}{\sqrt{2}(4\pi)^4} m_a m_b v_\xi y_{ab} \sin 2\theta \\ &\times \left[ \frac{1}{M_{P_1}^2} \log^2 \left( \frac{M_W^2}{M_{P_1}^2} \right) - \frac{1}{M_{P_2}^2} \log^2 \left( \frac{M_W^2}{M_{P_2}^2} \right) \right],\end{aligned}\tag{4}$$

where  $m_{a,b}$  correspond to charged lepton masses;  $P_1$  and  $P_2$  are the mass eigenstates of doubly charged scalars with  $\theta$  representing their mixing angle and  $M_{P_i} > M_W$  assumed. Note that the neutrino masses are suppressed by the two-loop factor,  $SU(2)_L$  gauge coupling, charged lepton masses, mixing angle  $\theta$ , and VEV of  $\xi$ , respectively, without fine-tuning Yukawa

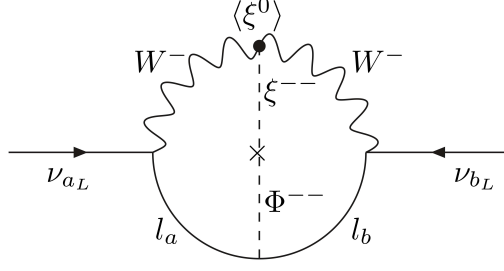


FIG. 1. Two-loop contributions to neutrino masses.

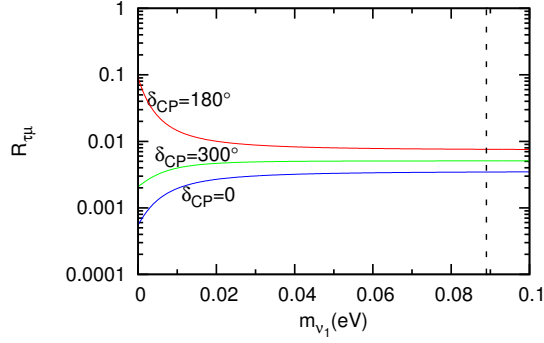


FIG. 2.  $R_{\tau\mu}$  versus  $m_{\nu_1}$ , where the blue, red and green curves correspond to  $\delta_{CP} = 0, 180^\circ$  and  $300^\circ$ , respectively, with the global  $\chi^2$  analysis in Ref. [17], while the dashed black line represents the lower bound set by cosmology.

couplings  $y_{ab}$ . The model predicts the neutrino mass spectrum to be a normal hierarchy if one requires the perturbative bound on  $y_{ab}$ . Consequently, the neutrino mass matrix is given by

$$m_{\nu_{ab}} = U_{\text{PMNS}} m_{\nu_{diag}} U_{\text{PMNS}}^T \propto y_{ab}, \quad (5)$$

where  $m_{\nu_{diag}} = \text{diag}(m_{\nu_1}, \sqrt{m_{\nu_1}^2 + \Delta m_{sol}^2}, \sqrt{m_{\nu_1}^2 + \Delta m_{atm}^2})$  with  $m_{\nu_1}$  the lightest  $\nu$  mass. From the neutrino oscillation data, one is able to pin down the neutrino parameters via the leptonic processes governed by  $y_{ab}$ . For example, the ratio  $R_{\tau\mu} \equiv \frac{Br(\tau \rightarrow e\gamma)^*}{Br(\mu \rightarrow e\gamma)}$  [18] is related to the lightest neutrino mass,  $m_{\nu_1}$ , as illustrated in Fig. 2, with the use of the latest neutrino oscillation data [17]. In principle, many variants of this quantity can be also defined from the leptonic rare decays or same-sign dilepton decays of  $\Phi^{\pm\pm}$ . Hence, our model provides a complementary way to determine neutrino parameters. Due to the indirect couplings between  $\Phi^{\pm\pm}$  and  $W^\mp W^\mp$ , our model reveals a mechanism for possible large neutrinoless

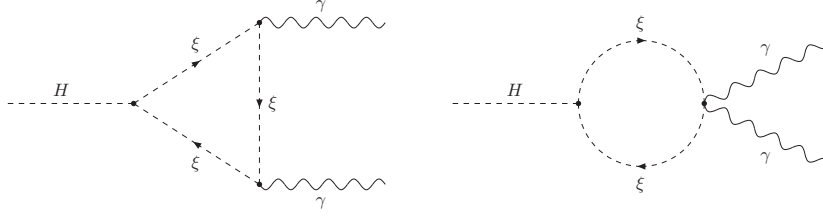


FIG. 3. Contribution to  $H \rightarrow \gamma\gamma$  from charged scalar exchanges in the loops.

double  $\beta$  decays which do not directly involve the small Majorana neutrino masses unlike the Zee-Babu model.

*LHC searches*— In our minimal model with  $\mathbf{n} = \mathbf{5}$ , apart from one SM-like Higgs scalar with  $m_H^2 = 2\lambda_H v^2$ , there are one pseudo-scalar, three singly charged, two doubly charged, and one triply charged scalars with masses all around a few hundred GeV, which all satisfy the current experimental bounds [16]. Since  $\xi$  does not directly interact with the SM fermions, the Higgs production is not expected to be modified. However, as promised, the decay rate of  $H \rightarrow \gamma\gamma$  receives extra contributions from the new charged scalars in the loops as shown in Fig. 3, given by

$$\begin{aligned} \Gamma(H \rightarrow \gamma\gamma) = & \frac{G_F \alpha^2 m_H^3}{128 \sqrt{2} \pi^3} \left| \sum_f N_f^c Q_f^2 A_{\frac{1}{2}}(\tau_f) + A_1(\tau_W) \right. \\ & \left. + \sum_{I_3} (I_3 + 1)^2 \frac{v}{2} \frac{\mu_s}{m_s^2} A_0(\tau_s) \right|^2, \end{aligned} \quad (6)$$

where the third term corresponds to those from the multi charged components of the dimension- $\mathbf{n}$  scalar multiplet  $\xi$  with their electric charges  $(I_3 + 1)$ ,  $I_3$  runs from  $(-\mathbf{n} + 3)/2$  to  $(\mathbf{n} + 1)/2$ , and  $\mu_s$  is the trilinear coupling to the SM Higgs. The amplitudes  $A_{0,\frac{1}{2},1}$  and the mass ratios  $\tau_{f,W,s}$  are defined in Ref. [19]. Here, for simplicity, we have taken the same trilinear coupling  $\mu_s$  and charged scalar mass  $m_s$ . The new contributions from the multi charged scalars interfere constructively with that of the SM if  $\mu_s < 0$ . As an illustration, we plot the ratio of  $R_{\gamma\gamma} \equiv \Gamma(H \rightarrow \gamma\gamma)/\Gamma(H \rightarrow \gamma\gamma)_{SM}$  in Fig. 4 with a typical value of  $\mu_s = -100$  GeV. It is clear that the excess of the rate reported at the LHC can be easily explained by these multi charged scalars.

The scalar multiplet components can be pair produced via  $q\bar{q} \rightarrow \gamma/Z^* \rightarrow \xi\xi^*(\Phi^{++}\Phi^{--})$  or  $q\bar{q}' \rightarrow W^* \rightarrow \xi^{+++}\xi^{--}(\xi^{++}\xi^-, \xi^+\xi^0)$  at the LHC. Similar processes for the triplet Higgs productions have been investigated in the Type-II seesaw mechanism [20]. The pair pro-

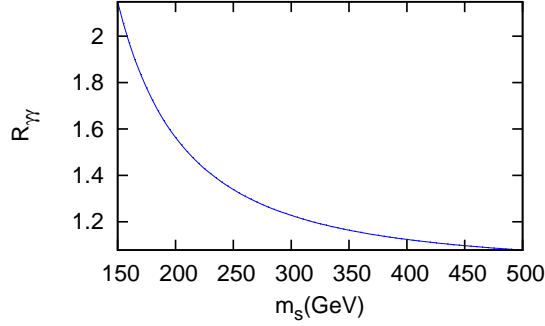


FIG. 4.  $R_{\gamma\gamma} \equiv \Gamma(H \rightarrow \gamma\gamma)/\Gamma(H \rightarrow \gamma\gamma)_{\text{SM}}$  as a function of the mass factor  $m_s$  from multi charged scalar states.

duction of  $\xi$  is estimated to have the order of  $10^{-3}\text{pb}$  with  $\sqrt{s} = 8 \text{ TeV}$  when  $m_\xi$  is around a few hundred GeV. For the doubly charged scalar, the lower mass limit has been set by CMS [21] and ATLAS [22] to be (382, 391, 395) and (409, 375, 398) GeV by assuming 100% branching ratio of  $(e^\pm e^\pm, e^\pm \mu^\pm, \mu^\pm \mu^\pm)$  final states, respectively. One noticeable feature of our model is that there are two doubly charged states,  $\xi^{\pm\pm}$  and  $\Phi^{\pm\pm}$ , in which the former is inert and only the latter couples directly to right-handed charged leptons. As a result, the lightest physical doubly charged scalar can be as small as  $\mathcal{O}(100) \text{ GeV}$  due to the mixing between  $\xi$  and  $\Phi$ , which would be observed at the LHC.

The signals of the triply charged component in the quintuplet scalar are distinct and its production rate is about two times larger than the doubly charged one according to their coupling strengths. The triply charged scalar  $\xi^{+++}$  decays into three W bosons with the width

$$\Gamma(\xi^{+++} \rightarrow 3W^+) \simeq \frac{3g_W^6 v_\xi^2 m_\xi^5}{512\pi^3 m_W^6}, \quad (7)$$

which subsequently go into the final states with six fermions, where we have ignored the phase space suppression. Note that in the degenerate limit, all W bosons can be on-shell. In addition to the multiplet  $\xi$ , the doubly charged scalars  $\Phi^{\pm\pm}$  also make the signals of our model more distinct at the LHC. If  $m_{\xi^{+++}} > m_{\Phi^{++}}$ , a channel opens up for the triply charged scalar  $\xi^{+++} \rightarrow \Phi^{++} + W^{*+}$  through the mixings, which can be distinguished from the  $3W^{*+}$  one as  $\Phi^{\pm\pm}$  only couple the same sign right-handed charged leptons with large center masses. By tagging the  $\tau$ s produced from  $\Phi$  or  $W$  one can pin down the chirality of  $\tau$  from its subsequent pion distribution spectrum [23].

The above results can be easily extended to other higher charged scalars. It is clear that these multi high charged scalars are well motivated and provide more interesting searching grounds at the LHC.

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